Controlling laser spectra in a phaseonium photonic crystal using maser

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We study the control of quantum resonances in photonic crystals with electromagnetically induced transparency driven by microwave field. In addition to the control laser, the intensity and phase of the maser can alter the transmission and reflection spectra in interesting ways, producing hyperfine resonances through the combined effects of multiple scattering in the superstructure.

1. INTRODUCTION

Optical properties of matter can be substantially modified and controlled by coherent excitations. Such systems can exhibit quantum coherence and interference effects, e.g., enhanced nonlinear effects[1], electromagnetically induced transparency (EIT)[2–4], giant Kerr nonlinearity[5–7], lasing without inversion[8–10], efficient nonlinear frequency conversions [11, 12], coherence Raman scattering enhancement via maximum coherence in atoms[13] and molecules[14], enhanced lasing[15, 16], coherent Raman umklappscattering[17], photodesorption[18] to name a few. Recently quantum coherence effects has been applied to a new domain on plasmonics and shown to benefit nanophotonics[19, 20].

Additional coherences in the presence of a microwave field, in addition to the optical fields have gained much attention in the past decade[21, 22]. For example, microwave field has been utilized to envision novel effects like electromagnetically induced transparency with amplification(EITA)[23] in superconducting qubits, subwavelength atom localization[24], circuit QED[25], simultaneous slow and fast light[26], gains without inversion in quantum systems with broken parities[27], quantum storage[28–30], sub-Raman generation[31].

On the other hand superstructures composed of exotic materials, such as superconductors in photonic crystals, show interesting optical features [32, 33]. Similarly photonic crystal composed of dielectric and controllable quantum coherence medium can provide new optical properties. Recently we analyzed one such superstructure, i.e. a superlattice composed of alternating layers dielectric and quantum coherence phaseonium [34] medium driven by a laser field [35]. Alternating layers with different optical properties can also be achieved by counter-propagating control laser [36] and temperature tuning of superconducting layers [37].

In this paper, we have extended the idea of controlling quantum resonances in photonic crystals with electromagnetically induced transparency by incorporating of a microwave coupling between the forbidden Raman transition ($|c\rangle \leftrightarrow |b\rangle$) as shown in Fig. 1, thus extending the regime of microwave excitation to a new domain. The optical fields couple the dipole allowed transitions ($|a\rangle \leftrightarrow |b\rangle$) and ($|a\rangle \leftrightarrow |c\rangle$). Similar scheme in atomic (^{87}Rb) vapor has already been experimentally demonstrated to exhibit control over EIT [38]. Here we show that the transmission and reflection profiles from single interface, double interface and superlattice structure show interesting features in the presence of the microwave field than its absence [32, 33]. In a closed loop both the amplitude and the phase of the microwave field can be applied to manipulate the dynamics of the system. The main results of the paper is shown in Figs. 3 and 4. where we have plotted the reflection, transmission and dispersion from the superlattice for different choices of phase and amplitude respectively of the microwave field.

It is worth mentioning here that the single atom formalism is applicable to dilute phaseonium particles such as doped transparent crystal and the model enables us to focus on the physics due to EIT without being drown in the many body formalism. The interplay between microwave resonance and multiple reflections and transmissions in multilayered structures produces novel resonant features. Our system looks like, but not exactly a distributed feedback laser (DFL)[39]. The mechanism behind EIT and quantum coherence with maser is linear and therefore it cannot be considered as a laser. The gain (R, T>1) is acquired from an external laser which controls the transmission and reflection spectra. This kind of coherent control is not available in DFL.

The paper is organized as follows, in section II we present our model and derive the equation of motion using the master equation approach. In section III we present our finding of the numerical simulations for reflection, transmission and dispersion profiles from single interface, double interface and superlattice in the presence of a microwave field coupled lower levels. We conclude our results in section IV where we have also discussed the application this approach.

2. MODEL AND EQUATION OF MOTION

Let us consider a three-level medium in Λ (Raman) configuration. Here $|a\rangle \leftrightarrow |c\rangle$ is excited by a control field

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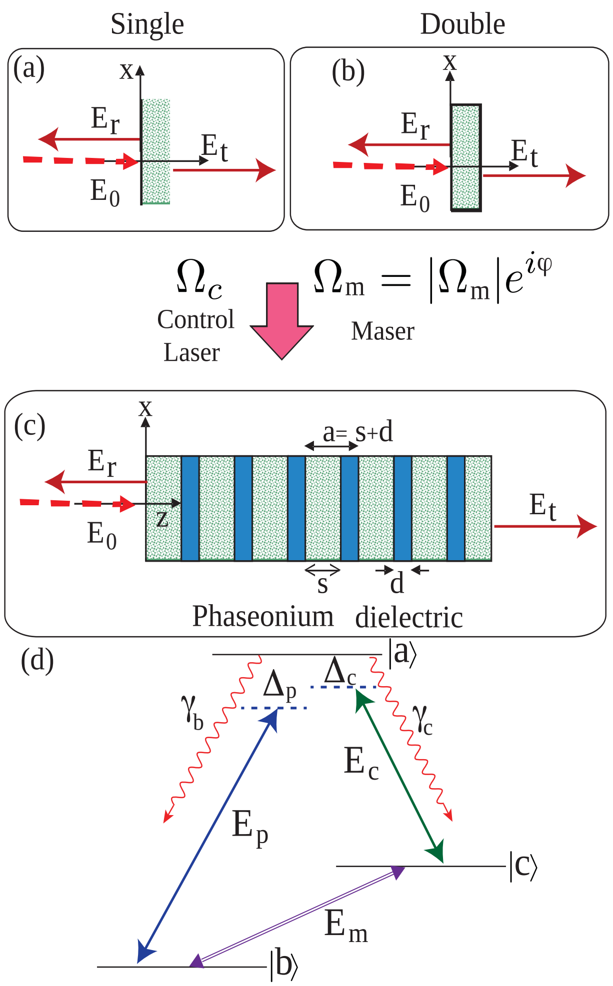

FIG. 1: Plot of the schematic of the superlattice with alternating layers composed of three-level scheme particles with EIT driven by control laser and a maser. Here E_0, E_r and E_t are the incident, reflected and transmitted probe field E_p.

E_c, while we probe the transition $|a\rangle \leftrightarrow |b\rangle$ with the field E_p. We have also coupled the transition $|c\rangle \leftrightarrow |b\rangle$ by a micro-wave field E_m. Let us define the (real) control, probe and microwave fields respectively as

$$E_c(z,t) = (1/2)\mathcal{E}_c(z,t)\exp[i(k_c z - \nu_c t)] + \text{c.c}, \qquad (1)$$

$$E_p(z,t) = (1/2)\mathcal{E}_p(z,t)\exp[i(k_p z - \nu_p t)] + \text{c.c}, \qquad (2)$$

$$E_m(z,t) = (1/2)\mathcal{E}_m(z)\exp[i(k_m z - \nu_m t)] + \text{c.c.} \qquad (3)$$

The atom-field interaction is given by

$$\begin{aligned}\mathscr{H} =& \hbar(\omega_a|a\rangle\langle a| + \omega_b|b\rangle\langle b| + \omega_c|c\rangle\langle c|) - \Big[\wp_{ab}E_p|a\rangle\langle b| \\ &+ \wp_{ac}E_c|a\rangle\langle c| + \wp_{cb}E_m|c\rangle\langle b| + \text{H.c}\Big].\end{aligned} \qquad (4)$$

Equation of motion for the density matrix ρ is given in

$$\begin{aligned}\frac{\partial\rho(z,t)}{\partial t} =& -\frac{i}{\hbar}[\mathscr{H},\rho] + \frac{\gamma_c}{2}\left([\sigma_c,\rho\sigma_c^\dagger] + [\sigma_c\rho,\sigma_c^\dagger]\right) \\ &+ \frac{\gamma_b}{2}\left([\sigma_b,\rho\sigma_b^\dagger] + [\sigma_b\rho,\sigma_b^\dagger]\right),\end{aligned} \qquad (5)$$

where the atomic lowering (σ_q) and raising ($\sigma_q^\dagger$) ($q = a,b,c$) operators are defined as $\sigma_c = |c\rangle\langle a|, \sigma_c^\dagger = |a\rangle\langle c|$; $\sigma_b = |b\rangle\langle a|, \sigma_b^\dagger = |a\rangle\langle b|$. To eliminate the fast oscillating terms we will use the following transformation.

$$\rho_{ab}(t) = \varrho_{ab}\exp[i(k_p z - \nu_p t)], \qquad (6)$$

$$\rho_{ac}(t) = \varrho_{ac}\exp[i(k_c z - \nu_c t)], \qquad (7)$$

$$\rho_{cb}(t) = \varrho_{cb}\exp[i(\delta k z - \delta\nu t)]. \qquad (8)$$

Here $\delta k = k_p - k_c$ and $\delta\nu = \nu_p - \nu_c$. The evolution of the coherences ϱ_{ij} takes the form

$$\begin{aligned}\frac{\partial\varrho_{ab}}{\partial t} =& -\Gamma_{ab}\varrho_{ab} - i\Omega_p(\varrho_{aa} - \varrho_{bb}) + i\Omega_c\varrho_{cb} \\ &- i\Omega_m\varrho_{ac}\exp[i(\Delta k_m z - \Delta\nu_m t)],\end{aligned} \qquad (9)$$

$$\begin{aligned}\frac{\partial\varrho_{ac}}{\partial t} =& -\Gamma_{ac}\varrho_{ac} - i\Omega_c(\varrho_{aa} - \varrho_{cc}) + i\Omega_p\varrho_{cb}^* \\ &- i\Omega_m^*\varrho_{ab}\exp[-i(\Delta k_m z - \Delta\nu_m t)],\end{aligned} \qquad (10)$$

$$\begin{aligned}\frac{\partial\varrho_{cb}}{\partial t} =& -\Gamma_{cb}\varrho_{cb} + i\Omega_c^*\varrho_{ab} - i\Omega_p\varrho_{ac}^* \\ &- i\Omega_m(\varrho_{cc} - \varrho_{bb})\exp[i(\Delta k_m z - \Delta\nu_m t)],\end{aligned} \qquad (11)$$

where we have defined the Rabi frequencies as $\Omega_p = \wp_{ab}\mathcal{E}_p/2\hbar, \Omega_c = \wp_{ac}\mathcal{E}_c/2\hbar, \Omega_m = \wp_{cb}\mathcal{E}_m/2\hbar$ and the detunings $\Delta k_m = k_m - \delta k$, $\Delta\nu_m = \nu_m - \delta\nu$. The decoherence Γ_{ij} are given as

$$\Gamma_{ab} = (\gamma_c + \gamma_b)/2 + i(\omega_{ab} - \nu_p) \qquad (12)$$

$$\Gamma_{ac} = (\gamma_c + \gamma_b)/2 + i(\omega_{ac} - \nu_c), \qquad (13)$$

$$\Gamma_{cb} = i(\omega_{cb} + \nu_c - \nu_p). \qquad (14)$$

Here we will consider the drive field is on resonance with the transition $|a\rangle \leftrightarrow |c\rangle$ i.e $\omega_{ac} = \nu_c$ while the probe and microwave field are detuned by Δ_p and Δ_m respectively. Analytical solution to Eqs. (9-11) can be found in different regimes and various limits [40]. Now assuming the drive field to be strong $|\Omega_c| \gg |\Omega_p|, |\Omega_m|$ and $\Delta_p = \Delta_m$, let us solve Eqs.(9-11) in steady state regime ($\dot\varrho_{ab} = \dot\varrho_{ac} = \dot\varrho_{cb} = 0$). We obtain,

$$\begin{aligned}\varrho_{ab} =& i\frac{\Gamma_{cb}\Omega_p\left(|\Omega_c|^2 n_{ac}/\Gamma_{ca}\Gamma_{cb} - n_{ab}\right)}{\Gamma_{ab}\Gamma_{cb} + |\Omega_c|^2} + \\ &\frac{\Omega_c\Omega_m\left(n_{cb} - \Gamma_{cb}n_{ac}/\Gamma_{ac}\right)\exp[i\Delta k_m z]}{\Gamma_{ab}\Gamma_{cb} + |\Omega_c|^2},\end{aligned} \qquad (15)$$

where $n_{ij} = \varrho_{ii} - \varrho_{jj}$. If we assume $\varrho_{bb} \sim 1, \varrho_{aa} = \varrho_{cc} \sim 0$, Eq.(15) reduces to

$$\bar{\varrho}_{ab} = i\frac{\Gamma_{cb}\Omega_p}{\Gamma_{ab}\Gamma_{cb} + |\Omega_c|^2} - \frac{\Omega_c\Omega_m \exp[i\Delta k_m z]}{\Gamma_{ab}\Gamma_{cb} + |\Omega_c|^2}. \quad (16)$$

Using the definition of complex susceptibility i.e $\chi = N|\wp_{ab}|^2\bar{\varrho}_{ab}/\hbar\epsilon_0\Omega_p$ and Eq.(16), we obtain,

$$\chi = \frac{N|\wp_{ab}|^2}{\hbar\epsilon_0}\left[i\frac{\Gamma_{cb}\left(|\Omega_c|^2 n_{ac}/\Gamma_{ca}\Gamma_{cb} - n_{ab}\right)}{(\Gamma_{ab}\Gamma_{cb} + |\Omega_c|^2)} + \frac{\Omega_c\Omega_m\left(n_{cb} - \Gamma_{cb}n_{ac}/\Gamma_{ac}\right)\exp[i\Delta k_m z]}{\Omega_p\left(\Gamma_{ab}\Gamma_{cb} + |\Omega_c|^2\right)}\right]. \quad (17)$$

Eq.(17) has two terms, the first term gives the well known $\Lambda-$scheme EIT, and the second term is the contribution from the microwave coupling of the transition $|c\rangle \leftrightarrow |b\rangle$. The net effect is determined by the interference of these two terms. The second term also determines the effect of relative phase of the optical and the microwave field on the probe field transmission. Absolute phase effects in radio-frequency(RF) domain has also been demonstrated by Jha *et.al.* [41, 42]. We use $k = \frac{\omega}{c}\sqrt{\epsilon_q(\omega)}$ with $\epsilon_q(\omega) = 1+\chi(\omega)$ to obtain the dispersion of the probe field inside the medium, ω versus Re(k) and Im(k), as shown in Fig.(2). Here we use ω, a more common notation for spectral frequency, to replaces ν_p the probe frequency in Eq. (17).

3. RESULTS

To understand the effect of the microwave coupling on the transmission and reflection we considered a single, double interface and superlattice. For single interface between incident medium and quantum coherence medium q we use the Fresnel relations

$$r_{iq}^{(p)} = \frac{\epsilon_q k_{iz} - \epsilon_i k_{qz}}{\epsilon_q k_{iz} + \epsilon_i k_{qz}}, \quad (18)$$

$$t_{iq}^{(p)} = \frac{2\sqrt{\epsilon_q\epsilon_i}k_{iz}}{\epsilon_q k_{iz} + \epsilon_i k_{qz}} \quad (19)$$

For double interface

$$r_{io}^{''} = \frac{r_{iq} + r_{qo}f^2}{1 + r_{iq}r_{qo}f^2}, \quad (20)$$

$$t_{io}^{''} = \frac{t_{iq} + t_{qo}f^2}{1 + r_{iq}r_{qo}f^2} \quad (21)$$

where $f = \exp(ik_{qz}d)$. The transmittance for single interface, $T = \left(k_{iz}|t_{iq}^2|\right)/\left(k_{qz}\right)$, and for the double interface, $T = |t_{io}^{''}|^2$. The one-dimensional photonic crystal (superlattice) shown in Fig. 1 is composed of a finite number of dielectric-phaseonium pairs. Multiple reflections and transmissions of a probe field E_0 through the multilayers give the reflected field E_r=r E_0 and the transmitted field E_r=t E_0 where the reflection and transmission coefficients

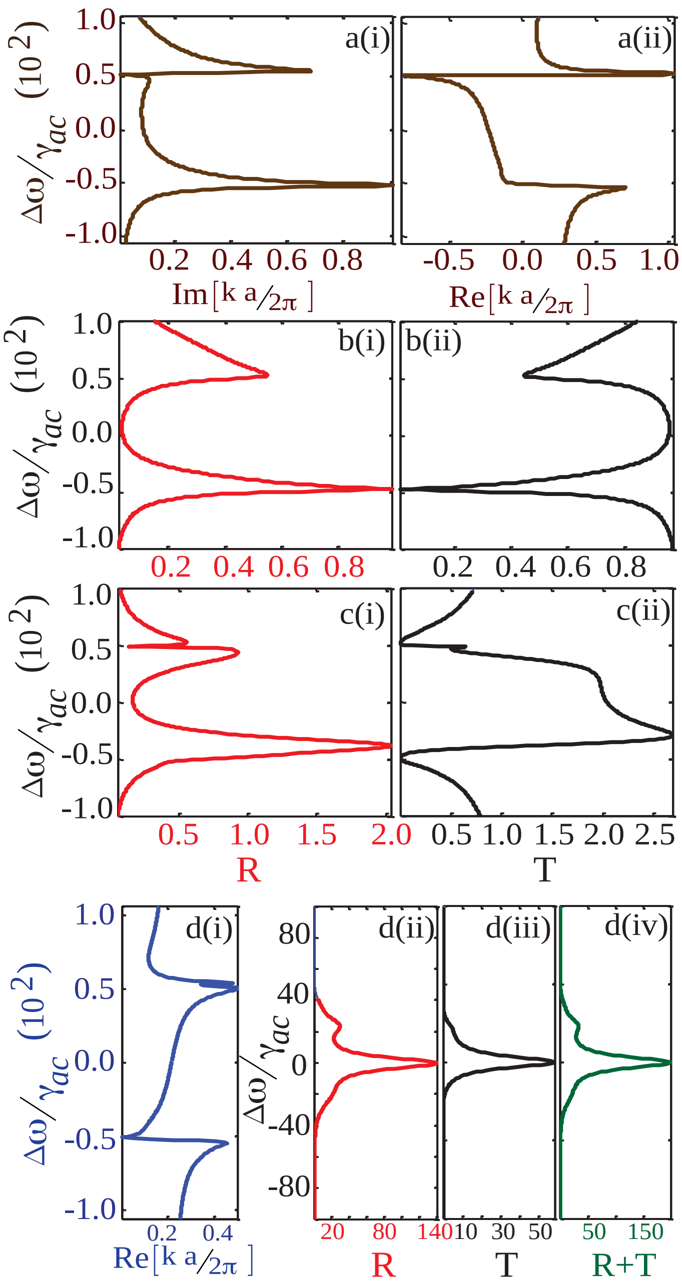

FIG. 2: Plot of absorption and dispersion from double interface a(i)-(ii) and dispersion from superlattice d(i). Reflection and transmission from single interface b(i)-(ii), , double interface c(i)-(ii) and superlattice d(ii)-(iii) in the presence of the maser field. For simulations we used $\Omega_c = 50\gamma_{ac}$, $\Omega_p = \gamma_{ac}$, and $\Omega_m = 0.02\gamma_{ac}$, $N = 10^{23}$, $\rho_{cc} = 0.0296$, $\rho_{aa} = 0.0155$. Here $a = 0.4\mu$m for the double interface (slab) and the superlattice and $d = s = 0.5a$ for the superlattice with $n = 8$.

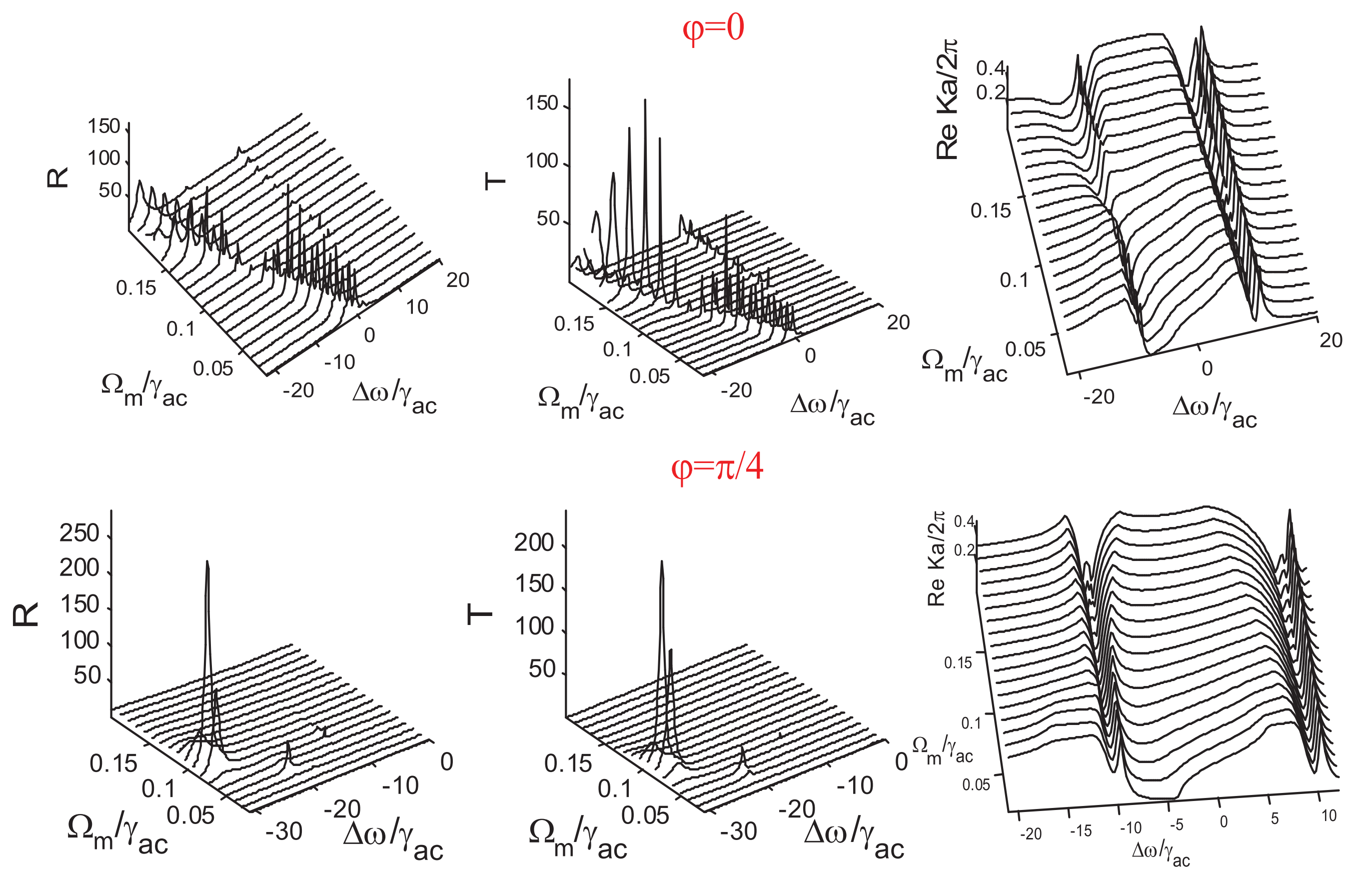

FIG. 3: Plot of the reflection, transmission and dispersion versus the Rabi frequency of the maser for a) $\varphi = 0$, b) $\varphi = \pi/4$, and c) $\varphi = \pi/2$. We use $\Omega_c = 10\gamma_{ac}$, $\Omega_p = 0.1\gamma_{ac}$, $\rho_{cc} = 0.0296$, $\rho_{aa} = 0.0155$ and other parameters are the same as in Fig. 2. Here we have defined φ as the phase of the microwave field.

are[43]

$$r = \frac{(M_{11} + M_{12}p_f)p_i - (M_{21} + M_{22}p_f)}{(M_{11} + M_{12}p_f)p_i + (M_{21} + M_{22}p_f)}, \tag{22}$$

$$t = \frac{s_{if}2p_i}{(M_{11} + M_{12}p_f)p_i + (M_{21} + M_{22}p_f)}, \tag{23}$$

where $p_j = \zeta k_{jz}/\omega c$, $\zeta = \sqrt{\varepsilon_0/\mu_0}$, $s_{if} = 1$ for TE-mode and $p_j = k_{jz}/\zeta k_j\sqrt{\varepsilon_j}$, $\sqrt{\epsilon_i/\epsilon_f}$ for TM-mode (ϵ_i and ϵ_f are the dielectric constants at the initial and final medium traversed by the probe.), M_{ij} are the components of the matrix

$$M = \begin{pmatrix} m'_{11}U_{N-1} - U_{N-2} & m'_{12}U_{N-1} \\ m'_{21}U_{N-1} & m'_{22}U_{N-1} - U_{N-2} \end{pmatrix} \tag{24}$$

m'_{ij} are components of 2×2 matrix $m' = (m_2 m_1)^{-1}$,

$$m_j = \begin{pmatrix} \cos\beta_j & i\frac{\sin\beta_j}{p_j} \\ ip_j\sin\beta_j & \cos\beta_j \end{pmatrix}, j = 1 \text{ or } 2 \tag{25}$$

with

$$\begin{aligned} U_N &= \frac{\sin(N+1)ka}{\sin ka}, \\ \beta_1 &= k_{1z}d,\ \beta_2 = k_{2z}s, \\ k_{jz} &= \frac{\nu}{c}\sqrt{\epsilon_j - \epsilon_i\sin^2\theta} \end{aligned} \tag{26}$$

and N is the number of phaseonium -dielectric layers. Also, note that $M = (m')^{N-1}$.

We note that the effective medium theory for susceptibility is only valid if the number of periods is very large and the thickness of the two layers are much smaller than the wavelength. For finite number of layers and the layers thickness in the order of wavelength, the present theory is exact and the effective medium theory is not a good approximation.

The results for the reflection, transmission for single interface, double interface and superlattice are shown in Fig. 2. We choose the control laser to be resonant to focus on results with clearer features. Figure 2b shows the usual two EIT peaks for single interface. Only for the double interface (Fig. 2c) and superlattice(Fig. 2d), R and T can exceed unity since the probe field is amplified by acquiring extra energy from the control field.

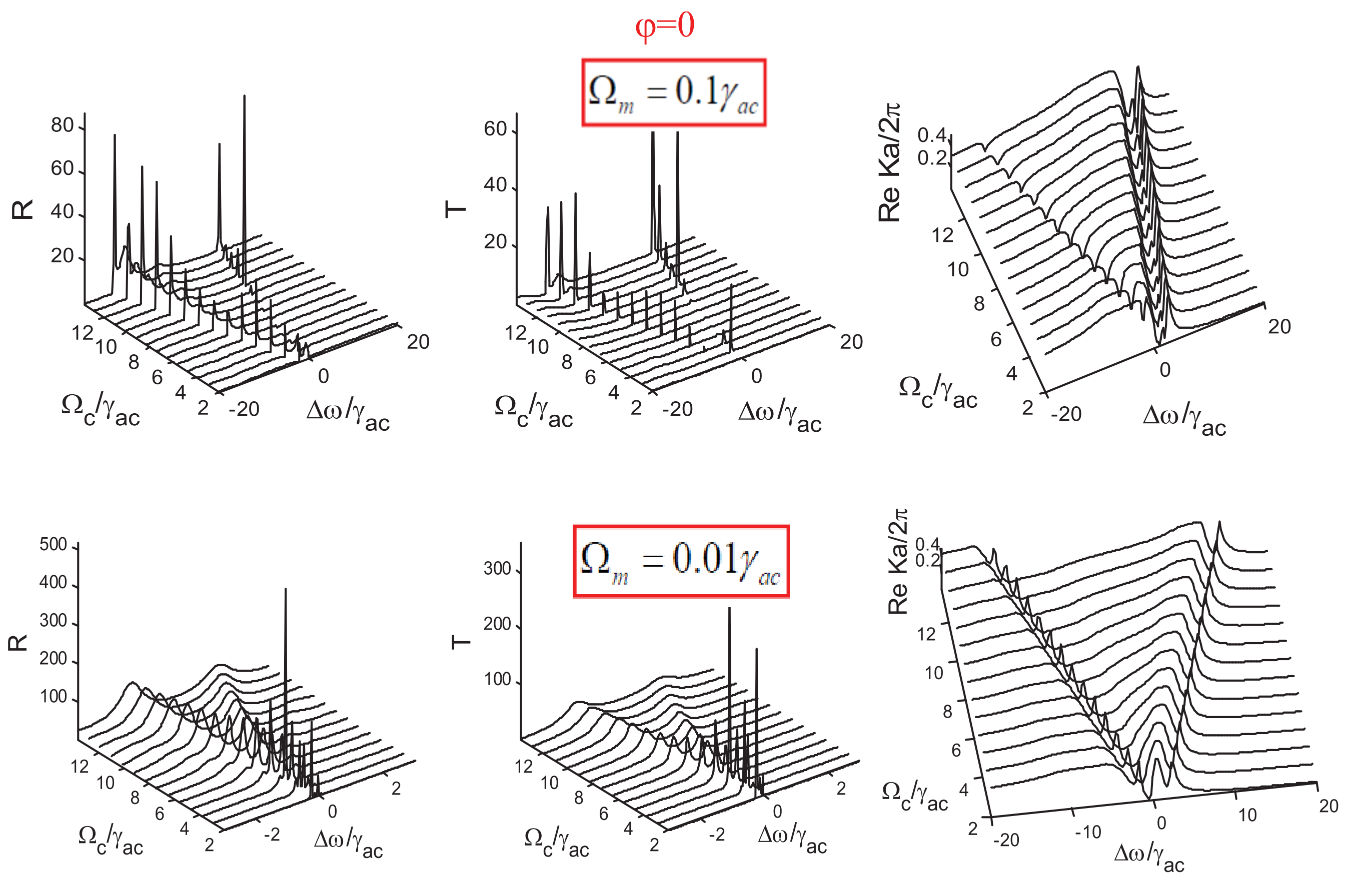

FIG. 4: Plot of the reflection, transmission and dispersion versus the Rabi frequency of the control field Ω_c for a) $\Omega_m = 0.1\gamma_{ac}$ and b) $\Omega_m = 0.01\gamma_{ac}$. We use $\Omega_p = 0.1\gamma_{ac}$, $\rho_{cc} = 0.0296$, $\rho_{aa} = 0.0155$ and other parameters are the same as in Fig. 2. The shift in the peaks becomes larger as Ω_m increases.

From Figs. 2b and c, for single and double interface we see that the reflection peaks appear around the two EIT resonances while the plateau of the transmission spectra lies between the resonances. The two EIT peaks corresponding to absorption give large reflection while the EIT window between the peaks gives high transmission if it does not fall in the photonic bandgap of the superlattice. However, there is a narrow transmission peak for double interface which corresponds to the dispersion and absorption in Fig. 2a, is the result of double reflection at the two interfaces. In general, multiple reflections lead to sharper reflection and transmission peaks. However, for the superlattice structure (see Fig. 2d) with the input energy from the control laser, there is a broad peak of both R and T are between the two EIT resonances. This is due to linear gain effect which has covered up the fine peaks from multiple reflections.

To study the effect of both the amplitude and the phase of microwave field we plotted (in Fig. 3) the dispersion, reflection and transmission as a function of Rabi frequency $|\Omega_m|$ for two choices of $\varphi = 0$ and $\pi/4$ where the distinct effect of the phase can be clearly seen from the results shown in Fig. 3. For $\varphi = 0$ the two (twin) EIT peaks are shifted to lower frequency as $|\Omega_m|$ increases to around $\Delta\omega/\gamma_{ac} = 0.1$. At this point, a small peak at higher frequency appears but the location of this peak does not shift with $|\Omega_m|$. Here, the EIT peaks becomes a single peak with the same size as the peak at higher frequency. As $|\Omega_m| > |\Omega_c|$ the separation and width of the twin peaks increase with $|\Omega_m|$. Such peaks do not appear in the absence of the microwave field. For $\varphi = \pi/4$, the Rabi frequency becomes complex and we find the tall and narrow peaks in R and T around $|\Omega_m| \simeq 0.1\gamma_{ac}$. It is important to mention here that the fine resonant features exist only if n_{ac} is finite such that additional resonances are contributed by the term $\Gamma_{cb}n_{ac}/\Gamma_{ac}$.

In Fig. 4, we have plotted the dispersion, reflection and transmission as a function of Rabi frequency Ω_c for three choices of $\Omega_m = 0.1\gamma_{ac}$ (upper) and $0.01\gamma_{ac}$ (lower). A sufficiently large microwave field can promote/enhance the separation between the twin EIT peaks in the transmission and reflection spectra as Ω_c increases. In particular the case $|\Omega_m| = 0.1\gamma_{ac} = |\Omega_p|$ transforms the lower frequency EIT peak into interesting double peaks, one narrow and one broad when Ω_c is sufficiently large.

4. CONCLUSION

In summary, we have investigated the effects of microwave coupling along the dipole forbidden (Raman) transition, on the quantum resonances in photonics crystal with electromagnetic transparency. We envision new applications based on our study. For example, the ultra-narrow transmission peak could be useful in high-precision spectroscopy as well as ultraslow light buffer or optical memory in all-photonic circuits. The narrow peak with high transmission in the multilayer structure with finite gain could be used for ultra narrow wavelength selection.

Recently controlling Raman and sub-Raman generation with microwave field has also been proposed[31]. The presence of gain in active structures has been used to compensate for absorption loss, promoting the practical use of quantum coherence in metamaterials and photonic crystals to a wider domain. Such ideas have given the birth of novel devices like nanolaser[44–46] and recently coherence effects have also been reported[19, 20] in such configurations. Incorporating microwave along with laser in such devices can bring additional degree of freedom for enhancing the gain and controlling resonances. In particular NV center in diamond, coupled to surface plasmon (both localized and propagating) are good candidate for the realization of microwave induced control in plasmonics[47]. We note that extension of this work with microwave field beyond the system investigated here into the hot field of plasmonics would be beneficial. One major challenge for coherent control in plasmonics is to use a drive field which does not affect (heat up) the nanostructure. In general the plasmon resonances have bandwidth ~100nm. Thus, microwave(maser) can be partially (if not completely) decoupled from the plasmon resonance and serve as an external control parameter. In general such approach also opens the door for nano-photonic devices to be merged with microwave telecommunication devices. However, for the experimental realization of the results presented here by the atomic physics/ quantum optics community, alkali metals like Rb,Na, K in a dielectric host are the feasible candidates.

5. ACKNOWLEDGEMENT

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